

Neutron Stars with Hyperons subject to Strong Magnetic Field

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Abstract

Neutron stars are one of the most exotic objects in the universe and a unique laboratory to study the nuclear matter above the nuclear saturation density. In this work, we study the equation of state of the nuclear matter within a relativistic model subjected to a strong magnetic field. We then apply this EoS to study and describe some of the physical characteristics of neutron star, especially the mass-radius relation and chemical compositions. To study the influence of a the magnetic field and the hyperons in the stellar interior, we consider altogether six solutions: three different values of magnetic field to obtain a weak, a moderate and a strong influence, and two configurations: a family of neutron stars formed only by protons, electrons and neutrons and a family formed by protons, electrons, neutrons, muons and hyperons. In all cases the particles that constitute the neutron star are in β equilibrium and zero total net charge. In the end, we see that although a magnetar can reaches $2.50M_{\odot}$, arises a natural explanation of why we do not know pulsars with masses above $2.0M_{\odot}$.

Keywords:

Neutron stars, pulsars, strong magnetic field, hyperons.

1 Introduction

Neutron stars are compact objects maintained by the equilibrium of gravity and the degenerescence pressure of the fermions together with a strong nuclear repulsion force due to the high density reached in their interior. Since we do not know yet the precise and detailed structure and composition of the inner core of a neutron star, many models have been used to describe it. In the literature we can find some standard ones: hadronic neutron stars, quark stars, strange stars and hybrid stars [1–3].

In the present work we study a hadronic neutron star constituted by nucleons and hyperons and subject to a strong magnetic field. The presence of hyperons is justifiable since the constituents of neutron stars are fermions. So, according to the Pauli Principle, as the baryon density increases, so do the Fermi momentum and the Fermi energy. Ultimately the Fermi energy exceeds the masses of the heavier baryons [1]. On the other hand, some strange objects like the soft gamma-ray repeaters and anomalous X-ray pulsars can be explained assuming that these objects are neutron stars subject to a strong magnetic fields on their surface. These objects are called magnetars [4]. Although the magnetic field of the magnetars do not exceed $10^{15}G$ in their surface, it is well-accepted in the literature that the magnetic field in the core of the neutron stars can reach values greater than

$10^{18}G$ [5,6]. Due to the large densities in the neutron star interior, we do not expect any significant influence of the magnetic field till it reaches values of the order of $10^{18}G$.

This paper is organized as follows: we make a review of the formalism of the non-linear Walecka model (NLWM) in the presence of a magnetic field. Then we present the numerical results showing how the presence of hyperons and a strong magnetic field affects the EoS and the chemical composition. Finally we study how these terms alter the macroscopic mass-radius relation of the neutron stars and compare our results with those found in literature.

2 The Formalism

The total Lagrangian is given by [7, 8]:

$$\mathcal{L} = \sum_b \mathcal{L}_b + \mathcal{L}_m + \sum_l \mathcal{L}_l + \mathcal{L}_B, \quad (1)$$

where b stands for the baryons, m for the mesons, l for the leptons, and B for the electromagnetic field itself. The sum in b can run over the eight lighter baryons and in l over the two lighter leptons. Explicitly, in the presence of a electromagnetic field the Lagrangian is:

$$\mathcal{L}_b = \bar{\Psi}_b [\gamma_u (i\partial^\mu - eA^\mu - g_{v,b}\omega^\mu - g_{\rho,b}I_{3b}\rho^\mu) - (M_b - g_{s,b}\sigma)] \Psi_b, \quad (2)$$

$$\begin{aligned} \mathcal{L}_m = & \frac{1}{2}\partial_\mu\sigma\partial^\mu\sigma - \frac{1}{2}m_s^2\sigma^2 + \frac{1}{2}m_v^2\omega_\mu\omega^\mu - \frac{1}{4}\Omega_{\mu\nu}\Omega^{\mu\nu} + \\ & + \frac{1}{2}m_\rho^2(\rho_\mu\rho^\mu) - \frac{1}{4}P_{\mu\nu}P^{\mu\nu} - \frac{1}{3!}\kappa\sigma^3 - \frac{1}{4!}\lambda\sigma^4, \end{aligned} \quad (3)$$

$$\mathcal{L}_l = \bar{\psi}_l [\gamma_u (i\partial^\mu - eA^\mu)] \psi_l, \quad (4)$$

$$\mathcal{L}_B = -\frac{1}{16\pi}F_{\mu\nu}F^{\mu\nu}, \quad (5)$$

where Ψ_b and ψ_l are the baryon and lepton Dirac fields, respectively. The baryon mass and isospin projection are denoted by M_b and I_{3b} respectively. The mass of leptons are m_l and the electric charge of the particles is given by e . The antisymmetric mesonic and electromagnetic field strength tensors are given by their usual expressions: $\Omega_{\mu\nu} = \partial_\mu\omega_\nu - \partial_\nu\omega_\mu$, $P_{\mu\nu} = \partial_\mu\rho_\nu - \partial_\nu\rho_\mu$ and $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. The γ_μ are the Dirac matrices [9]. The strong interaction couplings are denoted by g , and the meson masses by m , all with appropriate subscripts. The second subscript at the g constant is due to the distinctive coupling of hyperons with the mesons. In this work we assume that $g_{s,H} = 0.7g_{s,N}$; $g_{v,H} = 0.783g_{v,N}$ and $g_{\rho,H} = 0.783g_{\rho,N}$ [10], where H denotes hyperons and N nucleons. The hadronic part of the Lagrangian is the called NLWM. The leptons are included in the total Lagrangian density as a non-interacting Fermi gas in order to account the β equilibrium in the star.

To solve the equation of motion, we use the mean field approximation, where the meson fields are replaced by their expectation values, ie: $\sigma \rightarrow \langle\sigma\rangle = \sigma_0$, $\omega^\mu \rightarrow \delta_{0\mu}\langle\omega^\mu\rangle = \omega_0$ and $\rho^\mu \rightarrow \delta_{0\mu}\langle\rho^\mu\rangle = \rho_0$.

In this work we use a GM1 parametrization [11], which can describe the most important properties of nuclear matter and reproduce the macroscopic properties of the neutron stars consistent with those observed in nature. The GM1 parameters are:

Set	$(g_s/m_s)^2$	$(g_v/m_v)^2$	$(g_\rho/m_\rho)^2$	κ/M_N	λ
GM1	11.785 fm^2	7.148 fm^2	4.410 fm^2	0.005894	-0.006426

Table I: Values of GM1 parametrization.

This parametrization is fixed so that the incompressibility of nuclear matter $K = 300 \text{ MeV}$, and the nuclear saturation density $n_0 = 0.153 \text{ fm}^{-3}$. The masses of the baryon octet are $M_N = 939 \text{ MeV}$ (nucleons), $M_\Lambda = 1116 \text{ MeV}$, $M_\Sigma = 1193 \text{ MeV}$ and $M_\Xi = 1318 \text{ MeV}$. The meson masses are $m_s = 400 \text{ MeV}$, $m_v = 783 \text{ MeV}$ and $m_\rho = 770 \text{ MeV}$. And the masses of the leptons are $m_e = 0.511 \text{ MeV}$ and $m_\mu = 105.66 \text{ MeV}$. Applying the Euler-Lagrange in equation (1) in the absence of an electric field, the equation of motion in the mean field approximation for an arbitrary baryon becomes:

$$[\gamma_0(i\partial^0 - g_{v,b}\omega_0 - g_{\rho,b}I_{3b}\rho_0) - \gamma_j(i\partial^j - eA^j) - M_b^*]\Psi = 0, \quad (6)$$

where

$$M_b^* = M_b - g_{s,b}\sigma_0, \quad (7)$$

is the baryon effective mass.

For an uncharged particle eA^μ is always zero. The quantization rules are: $i\partial^0 = E$ and $i\partial^j = k^j$, where k^j is the momentum in j direction. Setting $E - g_{v,b}\omega_0 - g_{\rho,b}I_{3b}\rho_0 = E^*$, we have the following equation of motion written in a block matrix:

$$\begin{pmatrix} (E^* - M_b^*) & -\sigma \cdot \mathbf{k} \\ \sigma \cdot \mathbf{k} & -(E^* \rho_0 + M_b^*) \end{pmatrix} \begin{pmatrix} u_A \\ u_B \end{pmatrix} = 0. \quad (8)$$

This is an eigenvalue equation, which can be solved as the free Dirac equation for an effective mass and energy, whose solution is:

$$E = \sqrt{k^2 + M_b^{*2}} + g_{v,b}\omega_0 + g_{\rho,b}I_{3b}\rho_0 = \mu, \quad (9)$$

where μ is the chemical potential.

For a charged baryon, the Dirac equation assumes the following form:

$$\begin{pmatrix} (E^* - M_b^*) & -\sigma \cdot (\mathbf{k} - e\mathbf{A}) \\ \sigma \cdot (\mathbf{k} - e\mathbf{A}) & -(E^* - M_b^*) \end{pmatrix} \begin{pmatrix} u_A \\ u_B \end{pmatrix} = 0. \quad (10)$$

To produce a constant magnetic field in the z direction we choose: $A_2 = A_3 = 0$, $A_1 = -By$. The solution of this eigenvalue equation is well-know in the literature [5, 7, 12]:

$$E^* = \sqrt{M_b^{*2} + k_z^2 + 2\nu|e|B},$$

$$E = \sqrt{M_b^{*2} + k_z^2 + 2\nu|e|B} + g_{v,b}\omega_0 + g_{\rho,b}I_{3b}\rho_0 = \mu, \quad (11)$$

where the discrete parameter ν is called Landau level (LL) and μ is the chemical potential. The first LL , $\nu = 0$, is non degenerate and all the others are two-fold degenerate. For the leptons, since they don't feel the strong force:

$$E_l = \sqrt{m_l^2 + k_z^2 + 2\nu|e|B} = \mu_l. \quad (12)$$

And the expected values for the mesons are:

$$\omega_0 = \sum_{ub} \frac{g_{v,b}}{m_v^2} n^{ub} + \sum_{cb} \frac{g_{v,b}}{m_v^2} n^{bc} \quad (13)$$

$$\sigma_0 = \sum_{ub} \frac{g_{s,b}}{m_s^2} n_s^{ub} + \sum_{cb} \frac{g_{s,b}}{m_s^2} n_s^{bc} - \frac{1}{2} \frac{\kappa}{m_s^2} \sigma_0^2 - \frac{1}{6} \frac{\lambda}{m_s^2} \sigma_0^3, \quad (14)$$

$$\rho_0 = \sum_{ub} \frac{g_{\rho,b}}{m_\rho^2} n^{ub} I_{3b} + \sum_{cb} \frac{g_{\rho,b}}{m_\rho^2} n^{cb} I_{3b}, \quad (15)$$

where n^{cb} and n^{ub} are the number density of the “charged baryons” and “uncharged baryons” respectively [13, 14], and n_s^{cb} and n_s^{ub} are called scalar density for the charged and uncharged baryons [7]. In $T = 0$ ¹ they are given by:

$$dn^{ub} = \frac{8\pi k^2}{(2\pi)^3} \rightarrow n^{ub} = \int_0^{k_f} \frac{8\pi k^2}{(2\pi)^3} = \frac{k_f^3}{3\pi^2}, \quad (16)$$

$$dn^{cb} = \frac{|e|B}{(2\pi)^2} \eta(\nu) dk_z,$$

$$n^{cb} = \frac{|e|B}{(2\pi)^2} \sum_{\nu}^{\nu_{max}} \eta(\nu) \int_{-k_f}^{k_f} dk_z = \frac{|e|B}{2\pi^2} \sum_{\nu}^{\nu_{max}} \eta(\nu) k_f, \quad (17)$$

$$n_s^{ub} = \frac{1}{\pi^2} \int_0^{k_f} \frac{M_b^* k^2 dk}{\sqrt{M_b^{*2} + k^2}}, \quad (18)$$

$$n_s^{cb} = \frac{|e|B}{2\pi^2} \sum_{\nu}^{\nu_{max}} \eta(\nu) \int_0^{k_f} \frac{M_b^* dk_z}{\sqrt{M_b^{*2} + k_z^2 + 2\nu|e|B}}. \quad (19)$$

The summation in ν in the above expressions ends at ν_{max} , the largest value of ν for which the square of Fermi momenta of the particle is still positive and which corresponds to the closest integer from below defined by the ratio:

$$\nu_{max} \leq \frac{\mu^2 - M_b^{*2}}{2|e|B}, \quad \text{charged baryons} \quad (20)$$

$$\nu_{max} \leq \frac{\mu^2 - m_l^2}{2|e|B}. \quad \text{leptons} \quad (21)$$

Now we couple the equations imposing β equilibrium and zero total net charge:

$$\mu_{b_i} = \mu_n - e_i \mu_e, \quad \mu_e = \mu_\mu, \quad \sum_b e_b n_b + \sum_l e_l n_l = 0, \quad (22)$$

¹ This can justified since the Fermi temperature of the neutron stars is too high compared to its own temperature [15].

where μ_{b_i} and e_i are the chemical potential and electric charge of the i -th baryon, and μ_n , μ_e and μ_μ are the chemical potential of the neutron, electron and muon respectively.

The energy density of the neutron star is:

$$\epsilon = \sum_{ub} \epsilon_{ub} + \sum_{cb} \epsilon_{cb} + \sum_l \epsilon_l + \sum_m \epsilon_m + \frac{B^2}{8\pi}, \quad (23)$$

where the energy densities for the uncharged baryons, charged baryons, leptons and mesons have the following forms:

$$\epsilon_{ub} = \frac{1}{\pi^2} \int_0^{k_f} \sqrt{M_b^{*2} + k^2} k^2 dk, \quad (24)$$

$$\epsilon_{cb} = \frac{|e|B}{2\pi^2} \sum_{\nu}^{\nu_{max}} \eta(\nu) \int_0^{k_f} \sqrt{M_b^{*2} + k_z^2 + 2\nu|e|B} dk_z, \quad (25)$$

$$\epsilon_l = \frac{|e|B}{2\pi^2} \sum_{\nu}^{\nu_{max}} \eta(\nu) \int_0^{k_f} \sqrt{m_l^2 + k_z^2 + 2\nu|e|B} dk_z, \quad (26)$$

$$\epsilon_m = \frac{1}{2} m_s^2 \sigma_0^2 + \frac{1}{2} m_v^2 \omega_0^2 + \frac{1}{2} m_\rho^2 \rho_0^2 + \frac{1}{3!} \kappa \sigma_0^3 + \frac{1}{4!} \lambda \sigma_0^4. \quad (27)$$

To find the pressure, we use the second law of thermodynamics, that gives:

$$p = \sum_i \mu_i n_i - \epsilon + \frac{B^2}{8\pi}, \quad (28)$$

where the sum runs over all fermions. Note that the contribution from electromagnetic fields should be taken into account in the calculation of the energy density and the pressure.

2.1 TOV equations and the density-dependent magnetic field

The magnetic field of the surface of the magnetars are of order of $10^{15}G$, but can reach more than $10^{18}G$ in their cores. To reproduce this behaviour we use a density-dependent magnetic field given by [5, 12, 16]:

$$B(n) = B^{surf} + B_0 \left[1 - \exp \left\{ -\beta \left(\frac{n}{n_0} \right)^\alpha \right\} \right], \quad (29)$$

B^{surf} is the magnetic field on the surface of the neutron stars, taken as $10^{15}G$, n is the total number density, $n = \sum n^b$, B_0 is the constant magnetic field. The two free parameters β and α are chosen to reproduce a weak magnetic field below the nuclear saturation density, and a quickly growing when $n > n_0$, in such a way that $B(n) \sim B_0$ when $n > 6n_0$. To reproduce this behaviour, we have set $\beta = -5.76 \cdot 10^{-3}$ and $\alpha = 3$. Now B is replaced by $B(n)$ in the term $B^2/8\pi$ in our EoS.

To finish our analytical analysis, we write the TOV [17] equations:

$$\frac{dM}{dr} = 4\pi r^2 \epsilon(r), \quad (30)$$

$$\frac{dp}{dr} = -\frac{G\epsilon(r)M(r)}{r^2} \left[1 + \frac{p(r)}{\epsilon(r)} \right] \left[1 + \frac{4\pi p(r)r^3}{M(r)} \right] \left[1 - \frac{2GM(r)}{r} \right]^{-1}, \quad (31)$$

which are the differential equations for the structure of a static, spherically symmetric, relativistic star in hydrostatic equilibrium. The equation of states developed in this work are used as input for these equations. Since the distribution of the magnetic field is not spherical, we are aware that the TOV equations can only be used as an approximation in the present study.

3 Results and discussion

We consider two families of neutron stars: one containing just protons, electrons and neutrons, which we call “Atomic Stars” denoted by the letter A in the legends, and other containing protons, electrons, neutrons, muons and hyperons, which we call “Hyperonic Stars” denoted by the letter H in the legends. In the results we also include the crust of neutron star through the BPS EoS [18], but always taken into account the contribution of the magnetic field through the term $B(n)^2/8\pi$ in the EoS.

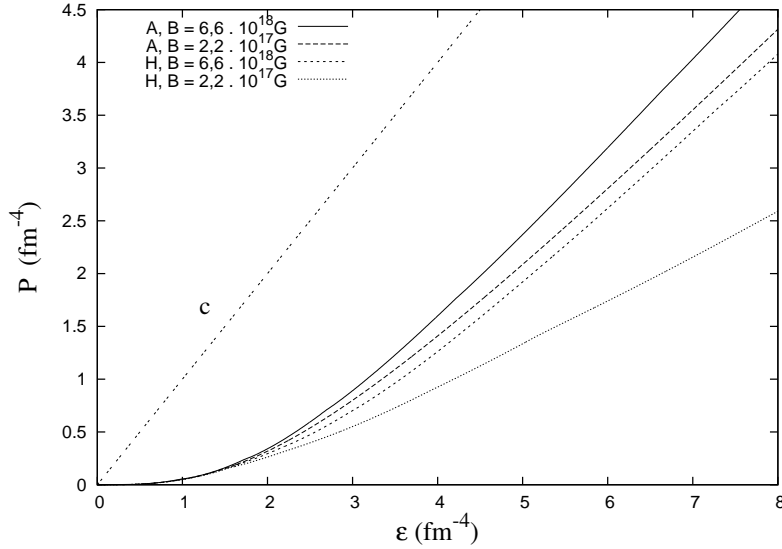


Fig. 1: EoS for two atomic and two hyperonic stars obtained with different values of the magnetic field. The straight line corresponds to the causal limit, for which $\epsilon = p$.

We choose three values for the magnetic field: $2.2 \cdot 10^{17}G$, $3.1 \cdot 10^{18}G$ and $6.6 \cdot 10^{18}G$, to produce a weak, a moderate and a strong influence. We also include here some theoretical and observational constrains. First, all our EoS are causal and obey the Le Chantelier principle, ie, the quantity $dp/d\epsilon$ lies between 0 and 1. We plot the numerical results of four EoS in fig. 1.

As we can see from fig. 1, the presence of hyperons softens the EoS more than the influence of the magnetic field can stiffen it. No matter how strong is the

magnetic field in the interior of the magnetar, the EoS of an atomic star is always stiffer than a hyperonic one. We also can see that all our EoS are causal (the “c” line is the causality limit).

From figs 1 and 3 we note that both the EoS and the fraction of particles $Y_i = n_i/n$ are not affected significantly by a magnetic field about $2.2 \cdot 10^{17}G$. The reason is that a field of this magnitude is too weak to contribute to the final EoS and to the fraction of particles. We also see that the magnetic field affects more the hyperonic stars than the atomic ones. As the hyperonic stars are softer than the atomic ones, they are therefore more sensitive to the presence of the magnetic field. Also the hyperonic stars are denser than the atomic stars, with a bigger central density n_c . Moreover, for a fixed value of density n , the pressure p of the hyperonic stars are smaller than the atomic ones. So, as the magnetic field couples to the number density through the equation (29), the contribution of the magnetic field is always greater in hyperonic than in atomic stars. We show this result plotting the pressure p in function of n in fig. 2.

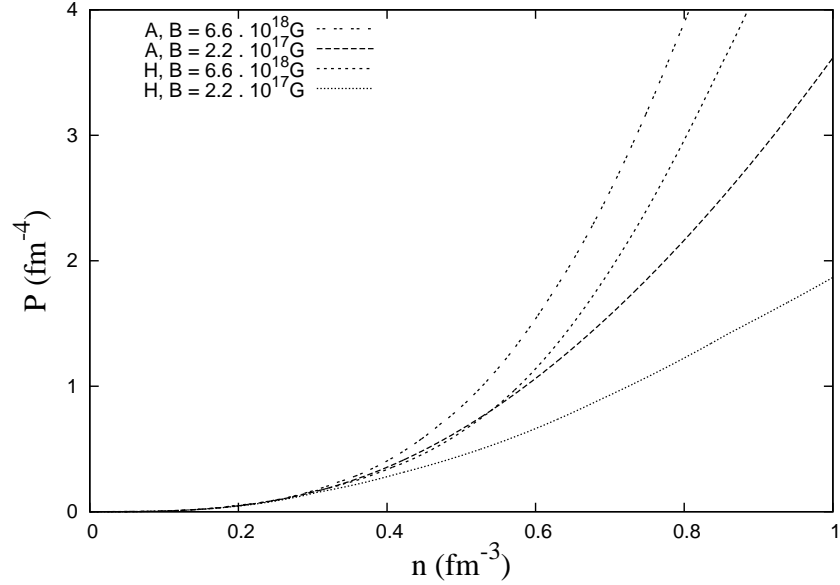


Fig. 2: Pressure as function of number density for four stars.

Figures 3, 4 and 5 show how the fraction of particles change with the increase of the magnetic field. We can see that for a considerable strength field, the appearance of charged particles is favoured at low densities due to their dependence on the magnetic field, as expected from eq. (17). Another change is the strange behaviour of the particles in the presence of a considerable magnetic field. For weak fields, the population of a kind of particle is always well-behaved, while for the strong ones many kinks appear.

The reason is that in the absence of a magnetic field the number density of a determined kind of particle grows smoothly with the momentum. When a magnetic field is present, there is also a dependence of the discrete LL . For a weak magnetic field a lot of Landau levels are occupied, but for a strong magnetic field, just a few of them are filled. So the orbit normal to the z direction is tightly quantized. This effect is more evident in the hyperonic star.

Fraction of particles Y_i as a function of number density n .

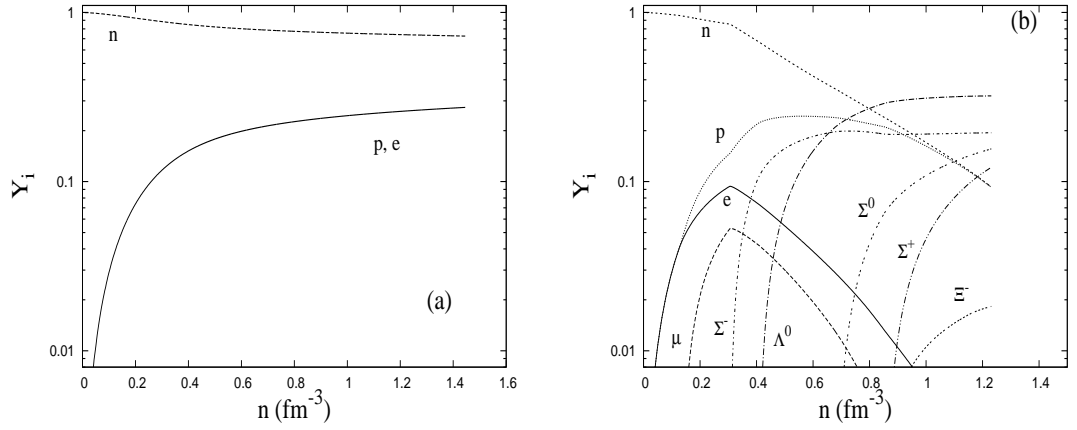


Fig. 3: (a) atomic stars and (b) hyperonic stars for a magnetic field of $2.2 \cdot 10^{17} \text{ G}$.

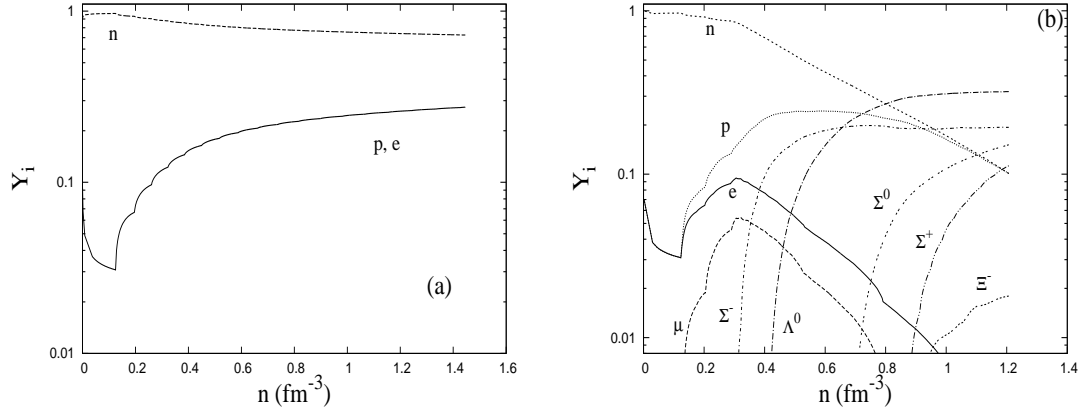


Fig. 4: (a) atomic stars and (b) hyperonic stars for a magnetic field of $3.1 \cdot 10^{18} \text{ G}$.

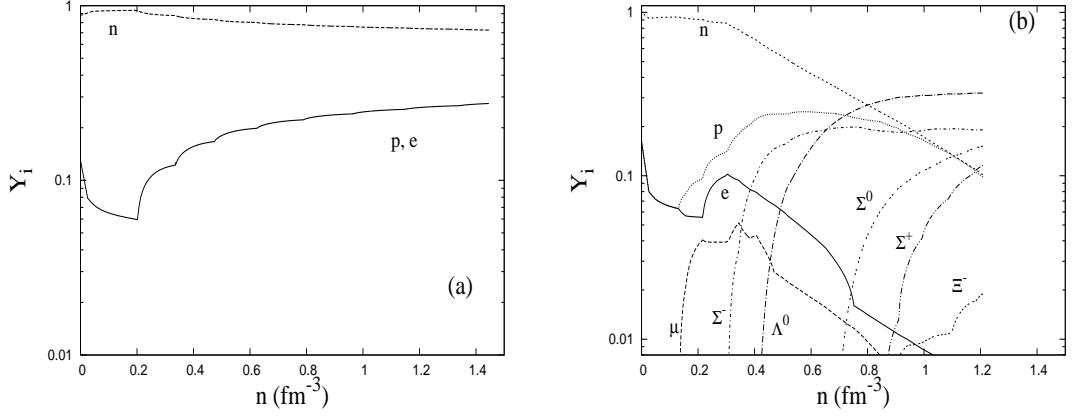


Fig. 5: (a) atomic stars and (b) hyperonic stars for a magnetic field of $6.6 \cdot 10^{18} \text{ G}$

Each nozzle in a slope of a determined particle indicates that the density is high enough to create another Landau level. For a high density there are so many Landau levels available that the distribution approaches to the continuous, while for a weak magnetic field, even in a low density there are several LL , so there is no significant difference with the case in the absence of the magnetic field. One can notice that for hyperonic stars, at densities of the order of 0.8 fm^{-3} the neutron is no longer the most important constituent. From this point, the Λ^0 hyperon dominates in the region of high densities.

It is hard to compare quantitatively our results with other ones existing in the literature due to the many different parametrizations. In the absence of hyperons, we see that our EoS is stiffer than those presented in [7, 8] since the authors of those references do not consider a density-dependent magnetic field. With relation to the fraction of particles Y_i , the fraction of protons in a low density region is more favoured in [7] than in our work. When hyperons are present, for a magnetic field of $2.2 \cdot 10^{17} \text{ G}$ there is no difference in both, the EoS and in Y_i when compared with the results for zero magnetic field [1]. When a strong magnetic field is applied, we see that our parametrization do not prevent the hyperon formation in a low density region as showed in [19]. Also, in our work, the influence of Landau quantization is much more evident even for a smaller value of the magnetic field.

In order to validate our EoS, we have to solve the TOV equations and check if the results agree with observational constraints. The star masses cannot exceed the maximum theoretical neutron star mass of 3.2 solar masses [20]. The EoS has to be able to predict the 1.97 solar masses neutron star [21] and to be in agreement with the redshift measurements (z) of two neutron stars. A redshift of $z = 0.35$ has been obtained from three different transitions of the spectra of the X-ray binary EXO0748-676 [22]. This redshift corresponds to $M/R = 0.15 M_\odot/\text{km}$. Another constraint on the mass-radius ratio comes from the observation of two absorption features in the source spectrum of the 1E 1207.4-5209 neutron star, with redshift from $z = 0.12$ to $z = 0.23$, which gives $M/R = 0.069 M_\odot/\text{km}$ to $M/R = 0.115 M_\odot/\text{km}$ [23]. Besides these constraints, the stars with central density above that of the maximum mass stars are mechanically unstable [1]. Due this fact,

the Ξ^0 hyperon is not present in the neutron star interior (the density required to create it is too high²).

Solving the TOV equations for the EoS we obtain the results presented in Table 2.

	Mass (M_\odot)	Radius (km)	n_c (fm^{-3})	B ($\times 10^{18}G$)
Max. Mass (A)	2.50	12.30	0.701	6.6
Max. Mass (A)	2.42	12.15	0.798	3.1
Max. Mass (A)	2.39	12.10	0.840	0.22
Max. Mass (H)	2.28	11.82	0.796	6.6
Max. Mass (H)	2.09	11.69	0.907	3.1
Max. Mass (H)	2.01	11.84	0.952	0.22

Table II Neutron stars properties computed from the six EoS used as input to the TOV equation. From the central density, we see that the hyperonic stars are denser than the atomic ones.

The fact that the EoS of hyperonic stars are always softer than the atomic ones reflects in the maximum mass of these stars. Further, from Table II we can see that the hyperonic stars are denser than the atomic ones, as said before. Ultimately, there is a curious behaviour of the radius of the maximum mass hyperonic stars. While in the atomic stars the radius grow with the magnetic field, in the hyperonic ones the radius has no visual dependence with it. We plot the TOV solutions in fig. 6 to compare our results:

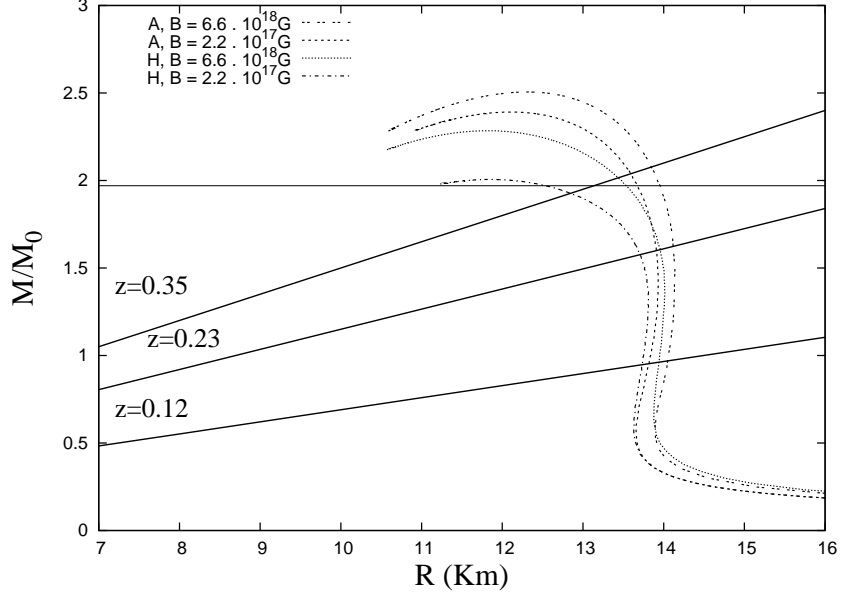


Fig. 6: Mass-radius relation for two atomic and two hyperonic stars with different values of magnetic field. The straight lines are the observational constraints

² indeed the Ξ^0 hyperon appears in an insignificant quantity $Y_{\Xi^0} = 10^{-5}$ at $1.2 fm^{-3}$.

In fig. 6 we can see that all our models are in agreement with the previously discussed constrains [20–23]. The inclined straight lines are the constrains of the measured redshift while the horizontal one is the $1.97 M_{\odot}$ pulsar. Every single dot in the curves is a possible neutron star. We can see also that a magnetar can reach a mass of $2.50 M_{\odot}$. However, as we do not know neither expected to discover any white dwarf with mass above $1.4 M_{\odot}$ due to the Chandrasekhar limit [2], the same behaviour appears to the neutron stars with mass above $2.0 M_{\odot}$ due to the Oppenheimer-Volkoff limit if we believe that hyperons and muons existing on their interior. The maximum possible mass is very close of the $1.97 M_{\odot}$ known neutron star. The discovery of a neutron star with mass above $2.0 M_{\odot}$ will imply that hyperons are absent or the star is a rare super magnetic neutron stars, as shown in fig. 6.

4 Conclusion

In this work we consider a hadronic neutron star composed by hyperons submitted to a strong magnetic field. We see that while the presence of hyperons reduce the maximum mass by softening the equation of state (EoS) [1, 24–27], the presence of a density-dependent magnetic field tends to increase the maximum mass stiffening the EoS [12, 16]. Our study shows also that although the influence of a strong magnetic field stiffens the EoS, a hyperonic magnetar still has a softer EoS compared with a common atomic stars. We also see how the magnetic field can change the chemical composition of neutron stars due to the Landau quantization and offer an explanation about the non-existence of neutron stars with mass above $1.97 M_{\odot}$. We conclude this paper by indicating that the effect of a strong magnetic field has to be taken into account in a description of magnetars, mainly if we believe that there are hyperons in their interior. In this case, the influence of the magnetic field can increase the mass by almost 15%.

Acknowledgments

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References

- [1] N. K. Glendenning: *Compact Stars - Nuclear physics, Particle physics, and General relativity*, Springer, New York - Second Edition - (2000)
- [2] S. L. Shapiro, S. A. Teukolsky: *Black Holes, White Dwarfs and Neutron Stars - The Physics of Compact objects*, John Wiley & Sons, New Jersey - (1983)
- [3] M. Camenzind *Compact Objects in Astrophysics: White Dwarfs, Neutron Stars and Black Holes* Springer-Verlag Berlin Heidelberg - (2007)
- [4] R. Duncan, C. Thompson: *Mon. Not. R. Astron. Soc.* - **275**, 255 - (1995)

- [5] S. Pal, D. Bandyopadhyay, S. Chakrabarty:
Phys. Rev. Lett. - **78**, 2898 - (1997)
- [6] S. Pal, D. Bandyopadhyay, S. Chakrabarty:
J. Phys. G: Nucl. Part. Phys. - **25**, L117 - (1999)
- [7] A. Rabhi, C. Providencia and J. Da Providencia:
J. Phys. G: Nucl. Part. Phys. - **35**, 125201 - (2008)
- [8] A. Broderick, M. Prakash, J. M. Lattimer:
Astrophys. J. - **537**, 351 - (2000).
- [9] D. Griffiths: *Introduction to Elementary Particles*,
WILEY-VCH, Weinheim - 2nd Edition - (2008)
- [10] M. G. Paoli, D. P. Menezes: *Eur. Phys. J. A* **46**, 413 - (2010)
- [11] N. K. Glendenning, S. A. Moszkowski:
Phys. Rev. Lett. - **67**, 2414 - (1991)
- [12] A. Rabhi et al:
J. Phys. G: Nucl. Part. Phys. - **36**, 115204 - (2009)
- [13] Kerson Huang: *Introduction to Statistical Physics*,
Taylor & Francis, London - (2001)
- [14] Q. Peng, H. Tong:
Mon. Not. R. Astron. Soc. **378**, 159 - (2007)
- [15] R. R. Silbar, S. Reddy:
Am. J. Phys. **72**, 7 - (2004)
- [16] D. P. Menezes et al: *Phys. Rev. C* - **80**, 065805 - (2009)
- [17] J. R. Oppenheimer, G. M. Volkoff:
Phys. Rev. - **33**, 374 - (1939)
- [18] G. Baym, C. Pethick, P. Sutherland:
Astrophys. J. - **170**, 299 - (1971)
- [19] A. Rabhi and C Providencia:
J. Phys. G: Nucl. Part. Phys. - **37**, 075102 - (2010)
- [20] C. E. Rhoades, R. Ruffini:
Phys. Rev. Lett. - **32**, 324 - (1974)
- [21] P. B. Demorest et al: *Nature.* - **467**, 1081 - (2010)
- [22] J. Cottam, F. Paerels, M. Mendez:
Nature - 420, 51 - (2002)
- [23] D. Sanwal et al.:
Astrophys. J. Lett. - 574, 61- (2002)
- [24] N. K. Glendenning: *Astrophys. J.* - **293**, 470 - (1985).

-
- [25] Z. X. Ma, Z. G. Dai, T. Lu :
Astronomy & Astrophysics - **366**, 532 - (2001)
- [26] J. L. Zdunik et al:
Astronomy & Astrophysics - **416**, 1013 - (2004)
- [27] H. Dapo, B. J. Schaefer, J. Wambach:
Phys. Rev. C - **81**, 035803 - (2010)